

Alternative Approach to Gaugino Condensation

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Abstract

We examine the mechanism of gaugino condensation in supersymmetric theories within a Nambu-Jona-Lasinio type approach. We investigate the effective Lagrangian description of higher energy theories that include some moduli fields in the gauge coupling constant. First we consider supersymmetric QCD with and without a mass term. We can find a phase transition in massless theory, but when we add a mass term, such a phase transition disappears. We also examine a model with a dilaton dependent coupling and find that it is very similar to supersymmetric QCD. Application of our method to supergravity is also examined.

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1 Introduction

There has recently been considerable attention focused on the study of supersymmetric models of elementary particle interactions. This is especially true in the context of grand unification theories, where remarkable studies have been done in the hope of solving the gauge hierarchy problem or unifying the gravitational interaction within the superstring formalism. Supersymmetric extension of the gravity(supergravity) seems necessary in introducing soft breaking terms and making the cosmological constant vanish simultaneously. In supergravity models, spontaneous breaking of local supersymmetry or super-Higgs mechanism may generate soft supersymmetry breaking terms that allow to fulfill such phenomenological requirements. However, the super-Higgs mechanism implies the existence of a supergravity breaking scale, intermediate between the Planck scale(M_p) and the weak scale(M_W). The intermediate scale is expected to be of $O(10^{13}\text{Gev})$. Here we expect that this intermediate scale is implemented by the mechanism of gaugino condensation in the hidden sector which couples to the visible sector by gravitational interactions. The effective action for gaugino condensation is well studied by many authors[1, 2]. The main purpose of this paper is to reproduce these results and seek for new features in these theories by means of the Nambu-Jona-Lasinio method[4, 5].

In section 2 and 3 we consider an intermediate scale effective Lagrangian for supersymmetric QCD with N_c colors and $N_f(< N_c - 1)$ flavors and also consider the effect of the dilaton dependent coupling constant. These theories have classical flat directions. If matter fields develop expectation value v on these directions, the original $SU(N_c)$ gauge symmetry is broken and the effective low energy Lagrangian has two parts. These are the kinetic terms for low energy pure supersymmetric Yang-Mills theory and the Nambu-Goldstone field, and their interaction term of order $O(1/v)$. Taking v to infinity, the resultant theory is a pure Yang-Mills theory without any interaction, so we tend to think that even in the existence of higher interaction terms we can extend the analysis of pure Yang-Mills theory. But this is merely a naive expectation so it seems important to analyze the theory from another point of view. From this standpoint, using the Nambu-Jona-Lasinio method we examine the effect of the higher terms that affect on gaugino condensation.

Section 4 includes the extension to supergravity.

2 Gaugino condensation in supersymmetric QCD

Gaugino condensation in supersymmetric gauge theories has been extensively studied by many authors both in global[1] and local[2] theories. In this section we examine the vacuum structures of Supersymmetric QCD(SQCD) theories with $N_f < N_c - 1$ by using the Nambu-Jona-Lasinio method. We follow ref.[1] in deriving the effective Lagrangian.

The results presented below nicely agree with the previous studies which are given by instanton or effective Lagrangian analysis.

Our starting point is a Lagrangian with a gauge group $SU(N_c)$ with N_f flavors of quarks. These superfields can be written with component fields as:

$$\begin{cases} Q^{ir} = \phi^{ir} + \theta^\alpha \psi_\alpha^{ir} + \theta^2 F^{ir} \\ \overline{Q}_{ir} = \overline{\phi}_{ir} + \overline{\theta}_\alpha \psi_{ir}^\alpha + \overline{\theta}^2 F_{ir} \end{cases} \quad (2.1)$$

The gauge fields $A_\mu^a (a = 1, \dots, N_c^2 - 1)$ are included in vector multiplets V^a accompanied by their super-partners, gauginos λ^a and auxiliary fields D^a . The total theory is given by

$$L = \frac{1}{4g^2} \int d^2\theta W^{\alpha a} W_\alpha^a + h.c. + \int d^4\theta [Q^+ e^V Q + \overline{Q} e^V \overline{Q}^+] \quad (2.2)$$

Classically, this theory has a global $U(N_f)_{Left} \times U(N_f)_{Right} \times U(1)_R$ symmetry. The $U(N_f)_{Left} \times U(N_f)_{Right}$ symmetry is just like that of ordinary QCD, corresponding to separate rotation of the Q and \overline{Q} fields. The symmetry $U(1)_R$ is a R-invariance, a symmetry under which the components of a given superfield transform differently. This corresponds to a rotation of the phases of the grassmannian variables θ^α ,

$$\begin{cases} \lambda & \rightarrow e^{i\alpha} \lambda \\ \psi & \rightarrow e^{i\alpha} \psi \\ \overline{\psi} & \rightarrow e^{i\alpha} \overline{\psi} \end{cases} \quad (2.3)$$

or

$$\begin{cases} W_\alpha(\theta) & \rightarrow e^{-i\alpha} W_\alpha(\theta e^{i\alpha}) \\ Q(\theta) & \rightarrow Q(\theta e^{i\alpha}) \\ \overline{Q}(\theta) & \rightarrow \overline{Q}(\theta e^{i\alpha}) \end{cases} \quad (2.4)$$

Just as in ordinary QCD, some of these symmetries are explicitly broken by anomalies. A simple computation shows that the following symmetry, which is a combination of the

ordinary chiral $U(1)$ and the $U(1)_R$ symmetry, is anomaly-free.

$$\begin{cases} W_\alpha(\theta) & \rightarrow e^{-i\alpha} W_\alpha(\theta e^{i\alpha}) \\ Q(\theta) & \rightarrow e^{i\alpha(N_c - N_f)/N_f} Q(\theta e^{i\alpha}) \\ \overline{Q}(\theta) & \rightarrow e^{i\alpha(N_c - N_f)/N_f} \overline{Q}(\theta e^{i\alpha}) \end{cases} \quad (2.5)$$

From now on, we call this non-anomalous global symmetry $U(1)_{R'}$.

Since this model has flat directions, it is reasonable to expect that Q and \overline{Q} may develop their vacuum expectation values along these directions. If $N_f < N - 1$, the gauge group is not completely broken. Moreover, we can see that instantons cannot generate a superpotential in this case, so considering another type of non-perturbative effects in this model seems important.

For simplicity, here we consider the case: $SU(N_c)$ gauge group is broken to $SU(N_c - N_f)$. The low-energy theory consists of two parts: Kinetic terms for the unbroken pure $SU(N_c - N_f)$ gauge interaction and one for the massless chiral field. In addition to these terms, we should include higher dimensional operators. A dimension-five operator, in general, is generated at one-loop level[3]. This can be obtained also from the renormalization of the effective coupling[1]:

$$L = \frac{1}{4g^2} \left[1 + \frac{g^2}{32\pi^2} N_f \ln \left(\frac{\phi}{\Lambda} \right) \right] W^\alpha W_\alpha \quad (2.6)$$

Of course, this term itself is not dimension five. Redefining the field as $\phi = \langle \phi \rangle + \phi'$, this term produces a dimension five operator, namely $\sim \frac{\phi'}{\langle \phi \rangle} W^2$. ϕ must be chosen to be invariant under all non-R symmetries. Detailed arguments on such an field dependence of coupling constant are given in ref.[1] and references therein. The non-anomalous R' -symmetry of the original theory must be realized in the effective low-energy Lagrangian by the shift induced by ϕ . That determines the R' -charge of ϕ to be $(N_c - N_f)/N_f$.

For simplicity, we consider a generalized form

$$L = \frac{1}{4} f(\phi) W^\alpha W_\alpha + h.c. + \phi^* \phi \quad (2.7)$$

where $f(\phi)$ is the field dependent coupling constant.

$$f(\phi) = \frac{1}{g_0^2} + \beta \log \left(\frac{\phi}{\Lambda} \right) \quad (2.8)$$

Here β is a constant chosen to realize the anomaly free (mixed) R'-symmetry of the original Lagrangian. In our case, we take $\beta = \frac{N_f}{32\pi^2}$. $\phi\phi^*$ in eq.(2.7) is not calculable and one may expect other complicated forms. Here we consider the simplest example for convenience.

The gauge group of the low energy theory is $SU(N_c - N_f)$. What we concern is the auxiliary part of this Lagrangian:

$$L_{AUX} = \frac{\beta g^2 \lambda \lambda}{v} F_\phi + h.c. + F_\phi^* F_\phi \quad (2.9)$$

(This term can be derived directly by 1-loop calculation.) We can simply assume that the cut-off scale of this effective Lagrangian is v . The factor of g^2 appears because we have rescaled gaugino fields to have canonical kinetic terms. The equation of motion for F_ϕ is:

$$\begin{aligned} \frac{\partial L}{\partial F_\phi} &= \frac{\beta g^2}{v} \lambda \lambda + F_\phi^* \\ &= 0 \end{aligned} \quad (2.10)$$

This equation means that $\langle \lambda \lambda \rangle$ is proportional to F_ϕ so we can think that $\langle \lambda \lambda \rangle$ is the order parameter for the supersymmetry breaking. Using the tadpole method[7] we can derive a gap equation directly from (2.10).

$$\begin{aligned} F_\phi^* \times \left(1 - 4G^2 \int \frac{d^4 p}{(2\pi)^4} \frac{1}{p^2 + m_\lambda^2} \right) &= 0 \\ \left\{ \begin{array}{l} G^2 = \frac{\beta^2 g^4}{v^2} (N_c - N_f) \\ m_\lambda^2 = \frac{|F_\phi|^2 g^4 \beta^2}{v^2} \end{array} \right. & \end{aligned} \quad (2.11)$$

Of course one can derive (2.11) by explicit calculation of 1-loop effective potential. Let us examine the solution. After integration we can rewrite it in a simple form.

$$\frac{4\pi^2}{G^2 \Lambda^2} = 1 - \left(\frac{m_\lambda^2}{\Lambda^2} \right) \ln \left(1 + \frac{\Lambda^2}{m_\lambda^2} \right) \quad (2.12)$$

(See also Fig.1a.) In the strong coupling region, this equation can have non-trivial solution. The explicit form of the potential is shown in Fig.1b. (Here we ignore the trivial solution $F_\phi = 0$ because such a condensation-vanishing solution does exist also in the effective (composite) Lagrangian analysis of pure Yang-Mills theory, but it is usually neglected.) Let us examine the behavior of this non-trivial solution. In pure supersymmetric Yang-Mills theories, gaugino condensation is observed even in the weak coupling region

because of the instanton calculation and Witten index argument that suggests the invariance of Witten index in the deformation of coupling constants[6]. If we believe that the characteristics of the low energy Lagrangian of massless SQCD is also similar to pure SYM, the weak coupling region should be lifted by gaugino condensation effect. On the other hand, if we believe that non-compactness of the moduli space is crucial and believe that gaugino condensation should vanish in the weak coupling region, we can think that the potential represented in Fig.1c is reliable and potential is flat in the weak coupling region. We cannot make definite answer to this question, but some suggestive arguments can be given by adding a small mass term to the field ϕ .

$$L_{mass}^{add} = \frac{1}{2}\epsilon\phi^2 \quad (2.13)$$

Existence of this term suggests that the moduli space is now compact. The resulting gap equation is drastically changed. We can naturally set F -components vanish, and the equation turns out to be a non-trivial equation for “ ϕ ”. Relevant terms are:

$$L_{AUX} = \left(\frac{\beta g^2}{v} \lambda \lambda + \epsilon \phi \right) F_\phi^* + h.c. + F_\phi^* F_\phi \quad (2.14)$$

The equation of motion for F_ϕ suggests that $\langle \lambda \lambda \rangle$ is now proportional to ϕ and no longer an order parameter for the supersymmetry breaking. The gap equation is given by:

$$\begin{aligned} \epsilon \phi \times \left(1 - 4G^2 \int \frac{d^4 p}{(2\pi)^4} \frac{1}{p^2 + m_\lambda^2} \right) = 0 \\ \begin{cases} G^2 = \frac{\beta^2 g^4}{v^2} (N_c - N_f) \\ m_\lambda^2 = \frac{\epsilon^2 g^4 \beta^2 |\phi|^2}{v^2} \end{cases} \end{aligned} \quad (2.15)$$

In general, this equation has a solution $m_\lambda = const.$ (see Fig.2a and 2b) which does *not* break supersymmetry(Fig.2c), and does not change Witten index for any(non-zero) value of ϵ and g_0 . In this case, the potential energy is always 0 for any value of g . Because the moduli space is compact in this case, it is reliable that there is no phase transition of gaugino condensation.

Now let us examine the limit $\epsilon \rightarrow 0$. In our model ϕ does not run away to infinity. If we take ϕ so large that the theory is weakly coupled, then non-trivial solution disappears and only the solution $\phi = 0$ is left. This contradicts the assumption, so we think ϕ is

finite and the potential is stabilized. Taking $\epsilon \rightarrow 0$, we can find a solution $4\pi^2/G^2\Lambda^2 = 1$ (see Fig.2d) i.e.:

$$\begin{aligned} G^2 &= \frac{\beta^2 g(\phi)^4}{\Lambda^2} (N_c - N_f) \\ &= \frac{4\pi^2}{\Lambda^2} \end{aligned} \quad (2.16)$$

This means that we can find a solution at $\phi = c\Lambda \exp(-\frac{1}{\beta g_0^2})$, here c is a constant ($c = \exp(\beta^2(N_c - N_f)/4\pi^2)$). Because we have fixed the symmetry breaking scale v and considered it as a cut-off scale for the low energy effective theory, we cannot find a runaway solution for $\langle \bar{Q}Q \rangle$ from this low energy Lagrangian.

3 Dilaton dependent coupling constant

In this section we mainly focus on the supersymmetric Yang-Mills theory in which the gauge coupling constant is dependent only on the dilaton field S . The Lagrangian is now written as:

$$L = \frac{1}{4}f(S)W^\alpha W_\alpha + h.c. - \Lambda^2 \log(S + \bar{S}) \quad (3.1)$$

Here we assume $Re f(S) = Re S \equiv \frac{1}{g_0^2}$. Relevant part of the Lagrangian is:

$$L_{AUX} = g^2 F_S \lambda \lambda + h.c. + \frac{F_S^* F_S}{(S + \bar{S})^2} \Lambda^2 \quad (3.2)$$

Here g means the renormalized coupling constant. Using the tadpole method we can find the following gap equation,

$$\begin{aligned} F_S^* &\times \left(\frac{\Lambda^2}{(S + \bar{S})^2} - 4G^2 \int \frac{d^4 p}{(2\pi)^4} \frac{1}{p^2 + m_\lambda^2} \right) = 0 \\ &\begin{cases} G^2 = g^4 N_c \\ m_\lambda^2 = g^4 |F_S|^2 \end{cases} \end{aligned} \quad (3.3)$$

which can be rewritten as:

$$\begin{aligned} F_S^* &\times \left(1 - 4G'^2 \int \frac{d^4 p}{(2\pi)^4} \frac{1}{p^2 + m_\lambda^2} \right) = 0 \\ &\begin{cases} G^2 = \frac{4}{\Lambda^2} \left(\frac{g}{g_0} \right)^4 N_c \\ m_\lambda^2 = g^4 |F_S|^2 \end{cases} \end{aligned} \quad (3.4)$$

This equation relies only on the parameter g/g_0 . This gap equation has the same characteristics of massless SQCD, so the dilaton potential is flat in the weak coupling region. The potential is almost the same as Fig.1c.

Then what would happen if we add a small mass term, for example, $L_m = \epsilon \Lambda^2 S^2/2$? Now the auxiliary part of the Lagrangian is:

$$L_{AUX} = \left[g^2 F_S \lambda \lambda + \epsilon \Lambda^2 S F_S + h.c. \right] + \frac{F_S^* F_S}{(S + \bar{S})^2} \Lambda^2 \quad (3.5)$$

This means that the dilaton field S is now becomes an order parameter for $\langle \lambda \lambda \rangle$.

The gap equation is now given by:

$$\epsilon S \times \left(1 - 4G^2 \int \frac{d^4 p}{(2\pi)^4} \frac{1}{p^2 + m_\lambda^2} \right) = 0$$

$$\left\{ \begin{array}{l} G^2 = \frac{g^4 (S + \bar{S})^2}{\Lambda^2} N_c \\ \quad = \frac{4}{\Lambda^2} \left(\frac{g}{g_0} \right)^4 N_c \\ m_\lambda^2 = \epsilon^2 |S|^2 g^4 (S + \bar{S})^4 \\ \quad = 16\epsilon^2 \left(\frac{g}{g_0} \right)^4 |S|^4 \end{array} \right. \quad (3.6)$$

This equation has a non-trivial solution at finite value once we fix g/g_0 .

Let us comment on the phenomenological models. If a small mass term is induced by the small spacetime curvature or some effects of higher theories, and assuming that our analyses are collect in such theories, we can expect that the analysis above may be used to analyze the dynamical dilaton potential and its runaway problem. It is important that we found a vacuum with finite S , without introducing multiple gauge groups.

4 Gaugino condensation in supergravity

In the standard superfield formalism of the locally supersymmetric action, we have:

$$S = \frac{-3}{\kappa^2} \int d^8 z E \exp \left(-\frac{1}{3} \kappa^2 K_0 \right) + \int d^8 z \mathcal{E} \left[W_0 + \frac{1}{4} f_0 \mathcal{W} \mathcal{W} \right] + h.c. \quad (4.1)$$

Here we set $\kappa^2 = 8\pi/M_p^2$. In the usual formalism of minimal supergravity, the Weyl rescaling is done in terms of component fields. However, in order to understand the anomalous quantum corrections to the classical action, we need a manifest supersymmetric

formalism, in which the Weyl rescaling is also supersymmetric. It is easy to see that the classical action(4.1) itself is not super-Weyl invariant. However, the lack of the super-Weyl invariance can be recovered with the help of a chiral superfield φ (Weyl compensator).

For the classical action (4.1), the Kähler function K_0 , the superpotential W_0 and the gauge coupling f_0 are modified[8]:

$$\begin{aligned} K_0 &\rightarrow K = K_0 - 6\kappa^{-2}\text{Re}log\varphi \\ W_0 &\rightarrow W = \varphi^3 W_0 \\ f_0 &\rightarrow f = f_0 + \xi log\varphi \end{aligned} \tag{4.2}$$

ξ is the constant chosen to cancel the super-Weyl anomaly. The super-Weyl transformations contain an R-symmetry in its imaginary part, so we can think that this is a natural extension of [5] in which a compensator for the R-symmetry played a crucial role.

Let us examine the simplest case. We include an auxiliary field H and set the form of W_0 and f_0 as:

$$\begin{aligned} W_0 &= \lambda H^3 \\ Ref_0 &= \frac{1}{g_0^2} \end{aligned} \tag{4.3}$$

and rescale the field φ as:

$$\tilde{\varphi} = H\varphi \tag{4.4}$$

where H is some auxiliary field. Finally we have:

$$\begin{aligned} K &= K_0 - 6\kappa^{-2}\text{Re}log\left(\frac{\tilde{\varphi}}{\Lambda}\right) \\ W &= \lambda\tilde{\varphi}^3 \\ f &= \frac{1}{g_0^2} + \xi log\left(\frac{\tilde{\varphi}}{\Lambda}\right) \end{aligned} \tag{4.5}$$

From the equation of motion for the auxiliary field of the super-Weyl compensator, we have the relation:

$$\lambda\tilde{\varphi}^3 - \frac{\xi}{6}g^2\lambda^\alpha\lambda_\alpha = 0 \tag{4.6}$$

And the tree level scalar potential is:

$$\begin{aligned} V_0 &= -3\kappa^2|W|^2 \\ &= -3\kappa^2\lambda^2|\tilde{\varphi}^3|^2 \end{aligned} \tag{4.7}$$

The equation of motion for the auxirialy field(4.6) suggests that eq.(4.7) can be interpreted as a four-fermion interaction of the gaugino:

$$-\frac{1}{12}\kappa^2 g^4 \xi^2 |\lambda^\alpha \lambda_\alpha|^2 \quad (4.8)$$

This four-fermion interaction becomes strong as $\frac{1}{g^2} = \text{Re}f$ reaches 0. The strong coupling point is:

$$\tilde{\varphi}_s = \Lambda e^{-\frac{1}{g_0^2}\xi} \quad (4.9)$$

Using the tadpole method we can have a gap equation:

$$\lambda \tilde{\varphi}^3 \times \left(1 - 4G^2 \int \frac{d^4 p}{(2\pi)^4} \frac{1}{p^2 + m_\lambda^2} \right) = 0$$

$$\begin{cases} G^2 = \frac{\xi^2 \kappa^2 g^4 N_c}{12} \\ m_\lambda^2 = \frac{\kappa^4 \xi^2 g^4 \lambda^2 |\tilde{\varphi}^3|^2}{4} \end{cases} \quad (4.10)$$

The solution for the gap equation(4.10) is given in Fig.3a. We can see that there is always a solution for non-zero gaugino condensation. (We have implicitly assumed that the vacuum expectation value for the auxiliary field for the Weyl compensator superfield is 0.) For a second example, we include the dilaton superfield S . Now f_0 is not a constant and depends on the field S :

$$f_0 = S \quad (4.11)$$

And the Kähler potential for the dilaton superfield is:

$$K_0 = -\kappa^{-2} \log(S + \bar{S}) \quad (4.12)$$

Here we should include the effect of the dilaton field in the scalar potential. The tree level scalar potential is:

$$V_0 = h_S (G^{-1})_S^S h^S - 3\kappa^2 |W|^2 \quad (4.13)$$

The auxirialy field for S is:

$$h_S = \kappa^2 \left[\frac{1}{2} \frac{W}{S + \bar{S}} + \frac{1}{4} f_S \lambda^\alpha \lambda_\alpha \right]$$

$$= \frac{\kappa^2}{4} W \frac{1 + 12 S_R \xi^{-1}}{S_R} \quad (4.14)$$

Here we set $G = K + \ln(\frac{1}{4}|W|^2)$ and $S_R = (S + \bar{S})/2$. The tree level potential can be given in a simple form

$$V_0 = \lambda^2 A |\tilde{\varphi}^3|^2 \quad (4.15)$$

where

$$A = \frac{1}{16} \kappa^2 \left[\left(1 + \frac{12S_R}{\xi} \right)^2 - 3 \right]. \quad (4.16)$$

In this case, the gap equation is given by:

$$\lambda \tilde{\varphi}^3 \times \left(1 - 4G^2 \int \frac{d^4 p}{(2\pi)^4} \frac{1}{p^2 + m_\lambda^2} \right) = 0$$

$$\begin{cases} G^2 = \frac{\lambda^2 \xi^2 g^4 N_c A}{36} \\ m_\lambda^2 = \frac{A^2 \xi^2 \lambda^2 |\tilde{\varphi}^3|^2}{36} \end{cases} \quad (4.17)$$

Now let us consider the difference between our result and ref.[5]. In ref.[5], the solution for the gap equation is estimated after fixing the coupling constant at g_c which is introduced by hand. It is true that the effective potential is singular at $\tilde{\varphi}_s$ (4.9), but without introducing the cut-off, we can find a solution for (4.17) at finite value.(see Fig 3a)

Another important point is the stability of the dilaton potential. It is suggested in [5] that the dilaton potential has a stable vacuum without introducing multiple gauge groups if we use the Nambu-Jona-Lasinio method. Related topics are also discussed in [9]. Can we explain this phenomenon along the line of the previous section? The scalar potential (4.15) contains a dilaton bilinear that can be interpreted as a mass term. As is shown in section 3, such a mass term would stabilize the dilaton potential.

5 Conclusion

We examined the formation of gaugino condensation in the hidden sector within a Nambu-Jona-Lasinio type approach. First we considered global supersymmetric gauge theories that have the $SU(N_c)$ gauge group and N_f matter fields. We can find the phase transition in massless SQCD, but in the massive theory, we cannot find such a phase transition. We can conclude that gaugino condensation is always non-zero in massive SQCD.

We also examined a model with a dilaton dependent coupling constant. The result is similar to SQCD. We found a stabilized dilaton potential when we add a small mass term. We also extended our analysis to supergravity.

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